HOW STUECKELBERG EXTENDS THE STANDARD MODEL AND THE MSSM

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Abelian vector bosons can get massive through the Stueckelberg mechanism without spontaneous symmetry breaking via condensation of Higgs scalar fields. This appears very naturally in models derived from string theory and supergravity. The simplest scenarios of this type consist of extensions of the Standard Model (SM) or the minimal supersymmetric standard model (MSSM) by an extra $U(1)_X$ gauge group with Stueckelberg type couplings. For the SM, the physical spectrum is extended by a massive neutral gauge boson Z′ only, while the extension of the MSSM contains a CP-even neutral scalar and two extra neutralinos. The new gauge boson Z′ can be very light compared to other models with $U(1)'$ extensions. Among the new features of the Stueckelberg extension of the MSSM, the most striking is the possibility of a new lightest supersymmetric particle (LSP) $\tilde{\chi}^0_{\rm St}$ which is mostly composed of Stueckelberg fermions. In this scenario the LSP of the MSSM $\tilde{\chi}^0_1$ is unstable and decays into $\tilde{\chi}^0_{\rm St}$. Such decays alter the signatures of supersymmetry and have impact on searches for supersymmetry in accelerator experiments. Further, with R-parity invariance, $\tilde{\chi}^0_{\rm St}$ is the new candidate for dark matter.

1. Stueckelberg mechanism for gauge boson masses

The Stueckelberg mechanism[1], as an alternative to spontaneous symmetry breaking in the Higgs effect[2], describes a way to make the naive Lagrangian of a massive abelian vector boson A_μ, such as

$$L_{\rm St} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{m^2}{2}A_\mu A^\mu \ ,$$

gauge invariant. To achieve this, one replaces $A_\mu \longrightarrow A_\mu + \frac{1}{m}\partial_\mu \sigma$, where σ is an axionic scalar that takes the role of the longitudinal mode of the massive vector, and defines the gauge transformation

$$\delta A_\mu = \partial_\mu \epsilon \ , \quad \delta\sigma = -m\epsilon \ .$$

The resulting Lagrangian is gauge invariant and renormalizable[3,4]. The physical spectrum contains just the massive vector field, and this mass growth occurs without the need for a charged scalar field developing a vacuum expectation value, without spontaneous symmetry breaking and accordingly without the need for a Higgs potential. In the above the mass parameter m is called a topological mass[4]. One can now go through the procedure of gauge fixing, so that A_μ and σ decouple in the final theory,

$$L_{\rm St} + L_{\rm gf} = -\frac{1}{4}F_{\mu\nu}F^{\mu\nu} - \frac{m^2}{2}A_\mu A^\mu - \frac{1}{2\xi}(\partial_\mu A^\mu)^2 - \frac{1}{2}\partial_\mu\sigma\partial^\mu\sigma - \xi\frac{m^2}{2}\sigma^2 \ .$$

A number of properties of this Lagrangian should be stressed. *i*) The vector A_μ has absorbed the real scalar σ in the process of getting a mass, with nothing left. *ii*) As the global subgroup of the gauge transformation, one can shift the scalar by a constant, $\delta\sigma = c$. This is a Peccei-Quinn like shift symmetry, and is the reason why we call σ an axionic pseudoscalar, which only appears with derivative couplings.[a] *iii*) Currently it appears possible to write such a gauge invariant Stueckelberg Lagrangian only for an abelian gauge symmetry, not for non-abelian gauge transformations[3]. However, as will become clearer from the string theoretic embedding of the Stueckelberg mechanism into D-brane models, the relevant $U(1)$ gauge group can become a subgroup of some non-abelian and simple grand unified gauge group in higher dimensions.

2. Stueckelberg couplings in string theory and supergravity

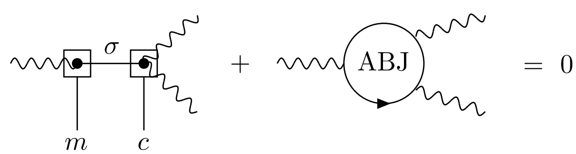
One immediate way to see that Stueckelberg couplings appear in dimensional reduction of supergravity from higher dimensions, and in particular string theory, is to consider the reduction of the ten-dimensional $N = 1$ supergravity coupled to supersymmetric Yang-Mills gauge fields, in the presence of internal gauge fluxes. The ten-dimensional kinetic term for the anti-symmetric 2-tensor B_{IJ} involves a coupling to the Yang-Mills Chern-Simons form, schematically $\partial_{[I}B_{JK]} + A_{[I}F_{JK]} + \cdots$, in proper units. Di-

[a] But σ does not necessarily have to couple to the QCD gauge fields in the usual topological term. In fact, we assume such couplings to be absent.

mensional reduction with a vacuum expectation value for the internal gauge field strength, $\langle F_{ij} \rangle \neq 0$, leads to

$$\partial_\mu B_{ij} + A_\mu F_{ij} + \cdots \;\sim\; \partial_\mu \sigma + m A_\mu \ ,$$

after identifying the internal components B_{ij} with the scalar σ and the value of the gauge field strength with the mass parameter m, which is indeed a topologiucal quantity, related to the Chern numbers of the gauge bundle. Thus A_μ and σ have a Stueckelberg coupling of the form $A_\mu \partial^\mu \sigma$. These couplings play an important role in the Green-Schwarz anomaly cancelation mechanism. In the effective four-dimenional theory, for instance abelian factors in the gauge group can have an anomalous matter spectrum, whose ABJ anomaly is canceled by Green-Schwarz type contributions. These involve the two terms $m A^\mu \partial_\mu \sigma + c\, \sigma F_{\mu\nu} \tilde{F}^{\mu\nu}$ in the Lagrangian.



As can be read from the left Feynman-diagram, the contribution to the anomalous 3-point function is proportional to the product of the two couplings, $m \cdot c$, while the mass parameter in the Stueckelberg coupling is only m. Therefore, any anomalous $U(1)$ will always get massive through the Stueckelberg mechanism, since $m \cdot c \neq 0$, but a non-anomalous $U(1)$ can do so as well, if $m \neq 0$, $c = 0$. Since we do not want to deal with anomalous gauge symmetries here, we shall always assume that $m \neq 0$, $c = 0$. The mass scale that determines m within models that derive from string theory can, at leading order, also be derived from dimensional reduction. It turns out to be proportional to the string or compactification scale in many cases[5], but can in principle also be independent[6].

The fact that an abelian gauge symmetry, anomalous or non-anomalous, may decouple from the low energy theory via Stueckelberg couplings was actually of great importance in the construction of D-brane models with gauge group and spectrum close to that of the SM[7]. Roughly speaking, these D-brane constructions start with a number of unitary gauge group factors, which are then usually broken to their special unitary subgroups via Stueckelberg couplings of all abelian factors, except the hypercharge,

$$U(3) \times U(2) \times U(1)^2 \quad \xrightarrow{\text{Stueckelberg}} \quad SU(3) \times SU(2)_L \times U(1)_Y$$

The mass matrix for the abelian gauge bosons is then block-diagonal, and only the SM survives. In order to ensure this pattern, one has to impose an extra condition on the Stueckelberg mass parameters, beyond the usual constraints that follow from the RR charge cancellation constraints, namely that the hypercharge gauge boson does not couple to any axionic scalar and remains massless[7]. In the language of these D-brane models, we will here relax this extra condition, and allow the hypercharge gauge boson to have Stueckelberg type couplings, and thus mix with other abelian gauge factors beyond the SM gauge group, which seems a very natural extension of the SM in this frame work.

3. Minimal Stueckelberg extension of the Standard Model

To keep things as simple as possible, and study the essence of the Stueckelberg effect in its minimal version[8,9], we consider an extension of the SM with only one extra abelian gauge factor $U(1)_X$. Along with the $U(1)_X$ we also allow a hidden sector with charged matter fields. Since the Stueckelberg mechanism cannot break the non-abelian $SU(2)_L$ in the process of electro-weak symmetry breaking, we cannot replace the usual Higgs mechanism, but only add the Stueckelberg masses for the abelian gauge bosons of hypercharge and $U(1)_X$ on top of the usual Higgs mechanism. Schematically, the couplings are thus given by

$$\begin{aligned} &SU(3) \times \overbrace{SU(2)_L \times U(1)_Y}^{\text{Higgs }\Phi} \times U(1)_X \\ &SU(3) \times SU(2)_L \times \underbrace{U(1)_Y \times U(1)_X}_{\text{Stueckelberg }\sigma} \ . \end{aligned}$$

The extra degrees of freedom beyond the SM spectrum are then the new gauge boson C_μ plus the axionic scalar σ which combine into the massive neutral gauge field Z′. Explicitily, we add the Lagrangian

$$L_{\rm St} = -\frac{1}{4} C_{\mu\nu} C^{\mu\nu} + g_X C_\mu J_X^\mu - \frac{1}{2}(\partial_\mu \sigma + M_1 C_\mu + M_2 B_\mu)^2$$

to the relevant part of the SM Lagrangian

$$\begin{aligned} L_{\rm SM} = -\frac{1}{4}{\rm tr}\, F_{\mu\nu}F^{\mu\nu} - \frac{1}{4}B_{\mu\nu}B^{\mu\nu} + g_2 A_\mu^a J_2^{a\mu} + g_Y B_\mu J_Y^\mu \\ -D_\mu \Phi^\dagger D^\mu \Phi - V(\Phi^\dagger \Phi) + \ \cdots \ . \end{aligned}$$

The scalars and vectors decouple after adding gauge fixing terms in a standard fashion. To keep the model as simple as possible, we impose the

following constraints on the charged matter spectrum: We assume that the fermions of the SM are neutral under $U(1)_X$, and vice versa we require that all the fields of the hidden sector, which are potentially charged under $U(1)_X$, be neutral under the SM gauge group.[b] Finally, one has to make sure that the hidden sector dynamics can really be ignored, i.e. that there is no spontaneous breaking of the $U(1)_X$ in the hidden sector.

4. Stueckelberg effects in the Standard Model

All effects of the minimal Stueckelberg extension on the SM Lagrangian can be summarized by the modified mass matrix of the, now three, neutral gauge bosons. In the basis (C_μ, B_μ, A^3_μ) for $U(1)_X$, hypercharge and the 3-component of iso-spin, The vector boson mass matrix is

$$\left[\begin{array}{c|cc} M_1^2 & M_1 M_2 & 0 \\ \hline M_1 M_2 & M_2^2 + \frac{1}{4} g_Y^2 v^2 & -\frac{1}{4} g_Y g_2 v^2 \\ 0 & -\frac{1}{4} g_Y g_2 v^2 & \frac{1}{4} g_2^2 v^2 \end{array}\right] \ .$$

The mass matrix above has one massless eigenstate, the photon with $M_\gamma^2 = 0$, and two massive ones with eigenvalues

$$M_{\mathrm{Z}}^2 = \frac{v^2}{4}(g_2^2 + g_Y^2) + \mathrm{O}(\delta) \ , \quad M_{\mathrm{Z}'}^2 = M^2 + \mathrm{O}(\delta) \ .$$

In the above we have introduced two parameters M and δ, defined by $M^2 = M_1^2 + M_2^2$ and $\delta = M_2/M_1$ to parametrize the Stueckelberg extension, which are a mass scale and a small coupling parameter. The bounds for these two, the values for which the Stueckelberg extension is still safely within experimentally allowed ranges, have been given in[8] as

$$M > [150\,\mathrm{GeV}] \ , \quad \delta < 0.01 \ .$$

For later use, we also define $M_0^2 = v^2(g_2^2 + g_Y^2)/4$. Roughly speaking, one may notice that all couplings which allow communication of SM fields to the Stueckelberg sector, are suppressed by δ. In models where an extra $U(1)'$ addition to the SM gauge group is broken by a Higgs condensate, the suppression factor comes from the propagator of the massive gauge boson, and is of the order of $M_{\mathrm{Z}}^2/M_{\mathrm{Z}'}^2 \sim 0.01$[10,11,12]. This demands that the mass

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[b]These conditions are very hard to satisfy in computable D-brane models, such as intersecting D-brane models with tori or toroidal orbifolds as compactification space. As a matter of principle, this should be an artifact of the too simple toroidal geometry and topology, and we expect that one can get around this problem in more general Calabi-Yau compactifications.

of the new gauge boson has to be in the range of TeV, much larger than the bound on M as above. For the Stueckelberg model, the suppression can easily be achieved by demanding δ be sufficiently small.[c]

A regime in the parameter space which would also be very interesting to investigate is the region $M_1,\ M_2 \to \infty,\ \delta = $ finite. This scaling limit corresponds to the expected behaviour in string theoretic models with a high mass scale, but basically arbitrary coupling δ. One may speculate that even though the Z′ becomes very heavy, and cannot be produced directly, small effects may remain observable. In essence, the Z′ does not decouple from the low energy theory, because the mass matrix is not diagonal, and this may have important consequences.

To proceed further, one can now diagonalize the vector boson mass matrix by an orthogonal matrix $O = O(\theta, \phi, \psi)$, a function of three angle variables $\{\theta, \phi, \psi\}$, with (see[8] for the details of the notation)

$$\tan(\phi) = \delta\ , \quad \tan(\theta) \ = \ \frac{g_Y}{g_2}\cos(\phi)\ , \quad \tan(\psi) = \frac{g_2(1 - M^2/M_{Z'}^2)}{g_Y \sin(\phi)\cos(\theta)}\ .$$

The bounds on M and δ translate into $\phi, \psi < 1^0$ and $\theta \sim \theta_W$, which is the electro-weak mixing angle. Inserting the mass eigenstates into the interaction Lagrangian

$$L_{\rm int} = g_2 A_\mu^a J_2^{a\mu} + g_Y B_\mu J_Y^\mu + g_X C_\mu J_X^\mu$$

one finds the interactions of the phyical vector fields. As the most striking example one gets for the photonic interactions

$$eA_\mu^\gamma J_{\rm em}^\mu = \frac{g_2 g_Y \cos(\phi)}{\sqrt{g_2^2 + g_Y^2\cos^2(\phi)}} A_\mu^\gamma \Big(J_Y^\mu + J_2^{3\mu} - \frac{g_X}{g_Y}\tan(\phi) J_X^\mu \Big)\ ,$$

which implies that the electric charge unit is now slightly redefined by

$$e = \frac{g_2 g_Y \cos(\phi)}{\sqrt{g_2^2 + g_Y^2\cos^2(\phi)}}$$

and that the charge unit of fields in the hidden sector, with charge under

[c]To date a real global fit of the experimental data within the Stueckelberg extended model has not been performed, and thus these statements cannot be made completely quantitative as yet.

$U(1)_X$, is irrational and very small.[d] Since all exotic couplings are suppressed by powers of δ, computable finite quantities that lead to definite predictions are best found in terms of ratios of such couplings. Examples are the branching ratios of the Z′, or the forward-backward asymmetry at the Z′ peak in resonant producation via e^+e^- collision[8]. One finds for example that the total width of the Z′ is extremely small, at least as long as the hidden sector charged fields are heavy enough, i.e.,

$$\Gamma(\mathrm{Z}' \to f\bar{f}) \;\sim\; \mathrm{O}(10)\,\mathrm{MeV}\ .$$

Therefore, Z′ would appear as a very sharp peak in e^+e^- annihilation and in other collider data.

5. Stueckelberg extension of the MSSM: StMSSM

The supersymmetrized version of the Stueckelberg coupling is related to the so-called linear multiplet formalism, see[13]. For the present minimal Stueckelberg extension of the MSSM, which we call StMSSM, it reads[9]

$$L_{\mathrm{St}} = \int d^2\theta d^2\bar{\theta}\ (M_1 C + M_2 B + S + \bar{S})^2\ .$$

Here S is the Stueckelberg chiral multiplet, and B, C are the abelian vector multiplets of hypercharge and $U(1)_X$. The degrees of freedom in components are given by $S = (\chi, \rho + i\sigma, F)$, $B = (B_\mu, \lambda_B, D_B)$, and $C = (C_\mu, \lambda_C, D_C)$, and the Lagrangian becomes

$$\begin{aligned} L_{\mathrm{St}} = -\frac{1}{2}(M_1 C_\mu + M_2 B_\mu + \partial_\mu\sigma)^2 - \frac{1}{2}(\partial_\mu\rho)^2 - \frac{i}{2}(\chi\sigma^\mu\partial_\mu\bar{\chi} - (\partial_\mu\chi)\sigma^\mu\bar{\chi}) \\ +\rho(M_1 D_C + M_2 D_B) + [\bar{\chi}(M_1\bar{\lambda}_C + M_2\bar{\lambda}_B) + \mathrm{h.c.}] + 2|F|^2\ . \end{aligned}$$

One may of course also add Fayet-Iliopoulos terms for each one of the two abelian factors, which would contribute to the scalar potential. As is usually done for the MSSM, we will however assume that their contributions to the breaking of supersymmetry are small compared to other sources, and therefore neglect these terms throughout, leaving a more complete analysis to future work. Eliminating the auxiliary fields F, D_B, D_C one thus finds corrections to the usual D-term potential of the MSSM through the coupling

[d] Note that a hidden sector with such an irrational electric charge is a possible consequence of the Stueckelberg scenario, which survives as a potentially observable fact in the stringy limit $M \to \infty$, $\delta =$ finite. While the mass of the gauge boson Z′ becomes very large, the hidden sector charged matter fields can be massless at the string scale, and thus be much lighter and in the end observable in experiment.

of the D-fields to ρ. To complete the action, we add the following soft supersymmetry breaking terms for the neutral gauginos, and the scalars $\{h_1, h_2, \rho\}$,

$$\begin{aligned} L_{\rm soft} = &-\frac{1}{2}\tilde{m}_\rho^2 \rho^2 - \frac{1}{2}\tilde{m}_1 \bar{\lambda}_B \lambda_B - \frac{1}{2}\tilde{m}_X \bar{\lambda}_C \lambda_C \\ &-\frac{1}{2} m_1^2 |h_1|^2 - \frac{1}{2} m_2^2 |h_2|^2 - m_3^2 (h_1 \cdot h_2 + \text{ h.c. }) \ . \end{aligned}$$

An important difference between the Stueckelberg Lagrangian and the Lagrangian that involves the Higgs mechanism is the fact that the chiral fermion χ is neutral under the gauge group. A scalar Higgs condensate would need to be charged under $U(1)_X$ to break the gauge symmetry spontaneously. But then, the standard coupling of the fermionic partner $\tilde{h}$ of the Higgs scalar in the form $g_Y B_\mu \tilde{h} \sigma^\mu \bar{\tilde{h}},\ g_X C_\mu \tilde{h} \sigma^\mu \bar{\tilde{h}}$ would imply a contribution ot the ABJ triangle anomaly, and a second Higgs multiplet of opposite charges would be needed to cancel the anomaly, just as in the MSSM. This is not the case for the Stueckelberg mechanism, where χ does not have such couplings, and no second chiral multiplet is needed.

Putting things together, the new effects in the StMSSM are in the scalar potential and the neutralino mass matrix, in addition to the mass matrix of the neutral gauge bosons, which is identical to that of the extended SM. The scalar potential with the three CP-even scalars $\rho,\ h_1,\ h_2$ is

$$\begin{aligned} V(h_1, h_2, \rho) &= \frac{1}{2}(M_1^2 + M_2^2 + \tilde{m}_\rho^2)\rho^2 + V_D^{\rm MSSM}(h_1, h_2) \\ &+\frac{1}{2}(m_1^2 - \rho g_Y M_2)|h_1|^2 + \frac{1}{2}(m_2^2 + \rho g_Y M_2)|h_2|^2 + m_3^2(h_1 \cdot h_2 + \text{ h.c. }) \ , \end{aligned}$$

where $V_D^{\rm MSSM}(h_1, h_2)$ is the usual D-term potential of the MSSM for the Higgs fields h_1, h_2 of MSSM. The new scalar ρ modifies the Higgs mass terms through its vacuum expectation value. Shifting $\rho \to v_\rho + \rho$ with $|g_Y M_2 v_\rho| < 10^{-4}\, M_Z^2$, one however finds that this induces only very tiny effects, as for instance in the electro-weak symmetry breaking constraint

$$\frac{1}{2} M_0^2 = \frac{m_1^2 - m_2^2 \tan^2(\beta)}{\tan^2(\beta) - 1} + \frac{g_Y M_2 v_\rho}{\cos(2\beta)} \ .$$

The CP-even scalar mass matrix in the basis (h_1, h_2, ρ) reads

$$\left[\begin{array}{cc|c} M_0^2 c_\beta^2 + m_A^2 s_\beta^2 & -(M_0^2 + m_A^2) s_\beta c_\beta & -t_\theta c_\beta M_{\rm W} M_2 \\ -(M_0^2 + m_A^2) s_\beta c_\beta & M_0^2 s_\beta^2 + m_A^2 c_\beta^2 & t_\theta s_\beta M_{\rm W} M_2 \\ \hline -t_\theta c_\beta M_{\rm W} M_2 & t_\theta s_\beta M_{\rm W} M_2 & M^2 + \tilde{m}_\rho^2 \end{array}\right] \ .$$

Going through the details, one again finds a very narrow resonance for the third mass eigenstate ρ_S in the $J=0^+$ channel, similar to the Z$'$, i.e.,

$$\Gamma(\rho_S \to t\bar{t}) \;\sim\; \mathrm{O}(10)\,\mathrm{MeV}\ .$$

Perhaps the most interesting sector of the StMSSM is the fermionic neutralino mass matrix, which now involves the usual two higgsinos and two gauginos of MSSM, plus the new gaugino of $U(1)_X$ and the chiral fermion χ of the Stueckelberg multiplet. In the basis $(\chi, \lambda_C, \lambda_B, \lambda_3, \tilde{h}_1, \tilde{h}_2)$ it reads

$$\left[\begin{array}{cc|cccc} 0 & M_1 & M_2 & 0 & 0 & 0 \\ M_1 & \tilde{m}_X & 0 & 0 & 0 & 0 \\ \hline M_2 & 0 & \tilde{m}_1 & 0 & -c_1 M_0 & c_2 M_0 \\ 0 & 0 & 0 & \tilde{m}_2 & c_3 M_0 & -c_4 M_0 \\ 0 & 0 & -c_1 M_0 & c_3 M_0 & 0 & -\mu \\ 0 & 0 & c_2 M_0 & -c_4 M_0 & -\mu & 0 \end{array}\right]$$

with abbreviations $c_1 = c_\beta s_\theta$, $c_2 = s_\beta s_\theta$, $c_3 = c_\beta c_\theta$, $c_4 = s_\beta s_\theta$, further s_θ, s_β, etc., standing for the sin and cos of θ_W and β, where $\tan(\beta) = \langle h_2\rangle / \langle h_1 \rangle$. It is convenient to number the eigenstates of the 4×4 MSSM mass matrix according to $m_{\tilde{\chi}_1^0} < m_{\tilde{\chi}_2^0} < m_{\tilde{\chi}_3^0} < m_{\tilde{\chi}_4^0}$, and the two new ones by

$$m_{\tilde{\chi}_5^0},\ m_{\tilde{\chi}_6^0} \;=\; \sqrt{M_1^2 + \frac{1}{4}\tilde{m}_X^2} \pm \frac{1}{2}\tilde{m}_X + \mathrm{O}(\delta)\ , \quad m_{\tilde{\chi}_5^0} \geq m_{\tilde{\chi}_6^0}\ .$$

The parameter space now easily allows the situation that the lighter one among the two new neutralinos becomes the lightest among all six. It therefore is the LSP, and with R-parity conservation, the natural dark matter candidate of the model. In this case, when $m_{\tilde{\chi}_6^0} < m_{\tilde{\chi}_1^0}$, we call $\tilde{\chi}_6^0$ the Stueckelino $\tilde{\chi}_{\mathrm{St}}^0$, and important modifications of the usual signatures for searches for supersymmetry in accelerator experiments[14] would follow. One would actually observe decay cascades (see also[15]), in which the lightest neutralino of the MSSM would further decay into the true LSP, the Stueckelino, via emitting fermion pairs,

$$\tilde{\chi}_1^0 \;\to\; l_i \bar{l}_i \tilde{\chi}_{\mathrm{St}}^0\ , \quad q_j \bar{q}_j \tilde{\chi}_{\mathrm{St}}^0\ , \quad \mathrm{Z}\tilde{\chi}_{\mathrm{St}}^0\ .$$

This would lead to decays of the form

$$\tilde{l}^- \;\to\; l^- + \tilde{\chi}_1^0 \;\to\; l^- + \begin{cases} l_i^- l_i^+ + \{\tilde{\chi}_{\mathrm{St}}^0\} \\ q_j \bar{q}_j + \{\tilde{\chi}_{\mathrm{St}}^0\} \end{cases}$$

for sleptons or

$$\tilde{\chi}_1^- \;\to\; l^- + \tilde{\chi}_1^0 + \bar{\nu} \;\to\; l^- + \begin{cases} l_i^- l_i^+ + \{\tilde{\chi}_{\mathrm{St}}^0 + \bar{\nu}\} \\ q_j \bar{q}_j + \{\tilde{\chi}_{\mathrm{St}}^0 + \bar{\nu}\} \end{cases}$$

for charginos, and similarly for squarks and gluinos. The analysis above shows that there will be multilepton final states in collider experiments which would be a characteristic signature for this of scenario.

6. Summary

We summarize now the most important features of the minimal Stueckelberg extension of the SM and of the MSSM:

(1) The Stueckelberg mechanism provides a gauge invariant, renormalizable method to generate masses for abelian gauge bosons with minimal extra residual scalar fields in the system. The mass of the massive vector boson is "topological" in nature.
(2) It naturally appears in many models that descend from string theory and higher-dimensional SUGRA.
(3) The Stueckelberg extension is very economical and distinct, even at the level of the degrees of freedom, compared to Higgs models with extra $U(1)'$ gauge factors.
(4) In the Stueckelberg extension of SM, only the vector boson sector is affected, as it introduces an extra Z$'$ boson which is typically a very narrow resonance. In addition, one has small exotic couplings of the photon and the Z with hidden matter, if present.
(5) In the Stueckelberg extension of the MSSM, the vector boson sector, the Higgs sector and the neutralino sectors are affected. The effect on the vector boson sector is identical to what one has in the Stueckelberg extension of the SM. In the neutral CP even Higgs sector one has mixing among the two CP-even Higgs of MSSM and a new CP-even Stueckelberg scalar field ρ. In the neutralino sector, one finds two more neutralinos which are mostly mixtures of neutral Stueckelberg fermions, in addition to the four neutralinos of MSSM.

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